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Multichannel calculation of the very narrow $D^*_{s0}(2317)$ and the very broad $D^*_0(2300\text{-}2400)$

George Rupp[1] and Eef van Beveren[2]

[1] Centro de Física das Interacções Fundamentais, Instituto Superior Técnico, Edifício Ciência, P-1049-001 Lisboa, Portugal
[2] Centro de Física Teórica, Departamento de Física, Universidade de Coimbra, P-3004-516 Coimbra, Portugal

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Abstract. The narrow $D^*_{s0}(2317)$ and broad $D^*_0(2300–2400)$ charmed scalar mesons and their radial excitations are described in a coupled-channel quark model that also reproduces the properties of the light scalar nonet. All two-meson channels containing ground-state pseudoscalars and vectors are included. The parameters are chosen fixed at published values, except for the overall coupling constant λ, which is fine-tuned to reproduce the $D^*_{s0}(2317)$ mass, and a damping constant α for subthreshold contributions. Variations of λ and $D^*_0(2300–2400)$ pole postions are studied for different α values. Calculated cross sections for S-wave DK and $D\pi$ scattering, as well as resonance pole positions, are given for the value of α that fits the light scalars. The thus predicted radially excited state $D^*_{s0}{}'(2850)$, with a width of about 50 MeV, seems to have been observed already.

PACS. 14.40.Lb Charmed mesons – 14.40.Ev Other strange mesons – 13.25.-k Hadronic decays of mesons – 12.39.Pn Potential models

The very narrow $D^*_{s0}(2317)$ charm-strange scalar meson first observed [1] three years ago has turned out to be the precursor of a series of new discoveries in hadron spectroscopy that have breathed new life into this field. The surprisingly low mass of the $D^*_{s0}(2317)$ itself has given rise to a flurry of theoretical work and speculations, mostly embracing non-standard quark configurations (see e.g. Ref. [2] for a list of references). Moreover, the very broad charm-nonstrange partner meson $D^*_0(2300–2400)$[1] discovered [4] shortly afterwards further added to the confusion, as its Breit-Wigner mass seems of the same order as the mass of the $D^*_{s0}(2317)$, and perhaps even larger [5]. However, the large width ($\simeq$ 260 MeV) of the $D^*_0(2300–2400)$ and the conflicting experimental mass determinations leave enough room for a possible reconciliation with the $D^*_{s0}(2317)$ mass.

In Ref. [6], we described the quasi-bound $D^*_{s0}(2317)$ and the $D^*_0(2300–2400)$ resonance as P-wave $c\bar{s}$ and $c\bar{n}$ ($n = u, d$) states, respectively, strongly coupled to the lowest S-wave two-meson channel, i.e., DK resp. $D\pi$. The framework of our calculation was a simple coupled-channel model previously used to fit the S-wave $K\pi$ phase shifts and predict the $K^*_0(800)$ (alias κ) meson [7]. As a result, both the quasi-bound $D^*_{s0}(2317)$ *below* the DK threshold and the very broad $D^*_0(2300–2400)$ resonance *above* the $D\pi$ threshold were roughly reproduced, though with a too low-lying $D^*_0(2300–2400)$ pole. Scaling arguments from flavour invariance later allowed to somewhat improve [2, 8, 9] our predictions, with no new parameters. In the present study (also see Ref. [10]), we ameliorate our coupled-channel description, by including all pseudoscalar-pseudoscalar (PP) and vector-vector (VV) channels, via a generalisation of our model that has recently been applied with success [11] to the whole light scalar nonet. The extension to PP and VV channels should allow us to make reliable predictions at least up to $\sim$3 GeV.

Inclusion of all (ground-state) PP and VV channels implies that we couple the scalar $c\bar{s}$ states to DK, $D_s\eta$, $D_s\eta'$ in S waves, and to D^*K^*, $D^*_s\phi$ in S as well as D waves, leading to a total number of 7 meson-meson channels. For the corresponding $c\bar{n}$ states, we need the coupling to $D\pi$, $D\eta$, $D\eta'$, D_sK in S waves, and to $D\rho$, $D\omega$, $D^*_sK^*$ in S and D waves, thus totalling 10 channels. For the (bare) confined $c\bar{q}$ states, an infinite harmonic-oscillator spectrum is taken, as in previous work. These bare states are then coupled, via the $^3P_0$ mechanism, to the two-meson channels, assuming that transitions only occur at a certain distance r_0. The resulting T-matrix can be solved in closed form (see Refs. [10, 11] for the formula).

An important and very difficult issue when dealing with coupled channels is how to treat subthreshold contributions, i.e., the effects of channels that are kinematically closed. Obviously, one cannot simply neglect channels as soon as the energy drops below threshold, which would be in gross violation of analyticity and even of common sense. However, it is also clear that, far below threshold, only a

[1] We adopt here the designation $D^*_0(2300–2400)$ instead of the official PDG nomenclature $D^*_0(2400)$ [3], to roughly indicate the two [4,5] observed mass values.

Table 1. Pole positions of $D^*_0(2300–2400)$ in MeV. Parameter λ fitted to $D^*_{s0}(2317)$ mass.

λ (GeV$^{-3/2}$)	r_0 (GeV^{-1})	Θ_{PS} (°)	α (GeV^{-2})	M_1M_2	$D^*_{s0}(2317)$	$D^*_0(2300–2400)$
2.491	3.2	-13.5	0.0	PP	2.317	$2186 - i109$
2.497	3.2	-17.3	0.0	PP	2.317	$2185 - i108$
1.468	3.2	-13.5	0.0	PP + VV	2.317	$2233 - i35.1$
1.469	3.2	-17.3	0.0	PP + VV	2.317	$2233 - i35.0$
2.700	3.2	-13.5	2.0	PP	2.317	$2165 - i111$
2.714	3.2	-17.3	2.0	PP	2.317	$2163 - i110$
2.137	3.2	-13.5	2.0	PP + VV	2.317	$2205 - i66.7$
2.144	3.2	-17.3	2.0	PP + VV	2.317	$2203 - i66.3$
2.854	**3.2**	**-13.5**	**4.0**	**PP**	2.317	$\mathbf{2149 - i111}$
2.868	3.2	-17.3	4.0	PP	2.317	$2147 - i110$
2.617	**3.2**	**-13.5**	**4.0**	**PP + VV**	2.317	$\mathbf{2174 - i96.4}$
2.629	3.2	-17.3	4.0	PP + VV	2.317	$2172 - i95.3$
2.988	3.2	-13.5	6.0	PP	2.317	$2135 - i108$
3.001	3.2	-17.3	6.0	PP	2.317	$2133 - i107$
2.901	3.2	-13.5	6.0	PP + VV	2.317	$2145 - i105$
2.913	3.2	-17.3	6.0	PP + VV	2.317	$2143 - i104$

non-perturbative field-theoretic treatment of the Dyson-Schwinger type might provide a rigorous description, since constituent masses are inexorably subject to major self-energy corrections in deeply bound systems. Evidently, a non-covariant approach like our coupled-channel Schrödinger equation cannot account for such effects, despite the use of relativistic kinematics, as particles are manifestly on-mass-shell. Suppression of closed channels due to wave functions, which is naturally included in our Schrödinger fomalism, empirically turns out to be insufficient in relativistic systems, as recently observed in the mentioned application of our model to the light scalars [11]. Therefore, we adopt here the same remedy as employed in the latter paper, and also in many multichannel data analyses, namely the use of subthreshold form factors. Thus, for closed channels we multiply the squares of the individual channel couplings that show up in our closed-form T-matrix expression by an exponential $\exp(\alpha k_i^2)$, where k_i is the relativistic channel momentum (with $\Re e\, k_i^2 < 0$) and α is a positive parameter, assumed to be universal. Clearly, this ansatz is not fully analytic either, namely on top of a threshold, but at least it is continuous there.

Now we proceed by fine-tuning the overall coupling constant λ [10] so as to reproduce the mass of the nowadays firmly established $D^*_{s0}(2317)$. We do this for a variety of situations in which not only the parameter α is chosen at different values, but also a comparison is made between calculations with only PP channels included, and with VV channels accounted for as well. Moreover, we also choose two different but both frequently quoted values for the pseudoscalar mixing angle θ_{PS}, which introduces slight variations in the predictions owing to the channels involving an η or η' meson. In each case, we determine the pole position of the ground-state scalar $c\bar{n}$ state. In Table 1, these pole positions are given together with the values of the parameters λ, α, θ_{PS}, and r_0 (fixed). From the table, we first of all observe that the dependence of the pole positions on the pseudoscalar mixing angle is indeed very feeble. Then, we note that the inclusion of the VV channels has a very significant effect on the $D^*_0(2300–2400)$ pole positions, especially on the imaginary parts, which nevertheless becomes smaller for increasing α. Also this can be easily understood, as all VV channels are highly virtual at these pole energies, so that large values of α lead to a strong suppression of these channels. Finally, all pole positions come out too low when compared to both experimental [4,5] $D^*_0(2300–2400)$ masses, even when noticing that the cross sections corresponding to these poles peak at somewhat higher masses. For instance, in the case $\alpha = 4.0$ GeV^{-2} (boldface in Table 1), which was the value used in Ref. [11] for all light scalars, our $D^*_0(2300–2400)$ cross-section peaks lie at 2.18 GeV (PP) and 2.19 GeV (PP+VV). However, some words of caution are due here. Besides the mentioned 100 MeV discrepancy between the two central experimental masses, it should be realised that one cannot just compare our predicted cross sections for elastic scattering with the Breit-Wigner fits of a very broad resonance observed in production processes, where other and more pronounced resonances like e.g. the $D^*_2(2460)$ show up as well. The resulting distributions for the $D^*_0(2300–2400)$ may be quite different (see e.g. FIG. 2 of Ref. [10]). Of course, pole positions must be the same in elastic scattering and production, but experiment does not extract any poles from the data. So better data on the $D^*_0(2300–2400)$ are definitely needed.

Focusing now our attention on the case $\alpha = 4.0$ GeV^{-2}, which fits the light scalars, we compute the elastic S-wave $D\pi$ and DK cross sections up to 3 GeV, which are then plotted in Fig. 1, both for the PP-only and the full PP+VV cases. Besides the large bump in $D\pi$ due to the $D^*_0(2300–2400)$, and the steeply falling cross section in DK owing to the $D^*_{s0}(2317)$ quasi-bound state, we ob-

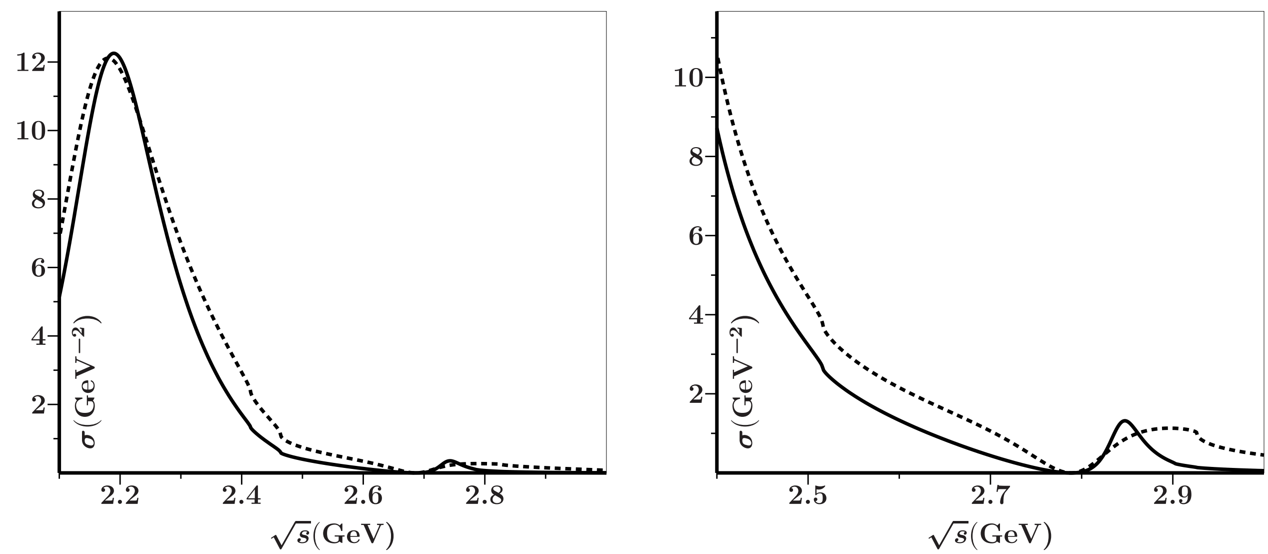

Fig. 1. S-wave $D\pi$ (left) and DK (right) cross sections. Dashed curves: PP channels only; full curves: PP+VV.

serve additional structures at higher energies, which are more pronounced and narrower when all channels are included. Concretely, there is a tiny bump in $D\pi$ at about 2.74 GeV, with a peak width of roughly 50 MeV. We find the corresponding pole at $2737 - i24.0$ MeV, which can be traced back, for vanishing λ, to the first radial excitation of the bare confinement spectrum at $E = 2823$ MeV. Moreover, there is a very broad pole as well, viz. at $2703 - i228$ MeV, which is connected to the confinement ground state, the D^*_0(2300–2400) being a continuum pole [2, 6]. In the DK case, there is a clear bump at about 2.85 GeV, with a peak width of again some 50 MeV. Besides a very broad continuum pole at $2779 - i233$ MeV, there is indeed a narrow pole at $2842 - i23.6$ MeV, originating from the bare confinement state at $E = 2925$ MeV. This resonance should thus correspond to the first radial excitation of the $D^*_{s0}(2317)$. Quite significantly, a new $c\bar{s}$ resonance denoted $D_{sJ}(2860)$, with a mass of $2856.6 \pm 1.5 \pm 5.0$ MeV and a width of $48 \pm 7 \pm 10$ MeV, was reported [12] by the BABAR collaboration [14] very shortly after the presentation of the present results. The observation of the DK decay mode and the non-observation of the D^*K mode, as reported by BABAR, are compatible with the radially excited scalar $c\bar{s}$ state predicted by us.

In the meantime, three other theoretical papers [15–17] on the $D_{sJ}(2860)$ have appeared. The first one argues in favour of a 3^- assignment, the second one supports a radially excited scalar as we do, and the third admits either option. So experimental confirmation of the $D_{sJ}(2860)$ is needed, as well as observation of another decay mode.

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